

HOW TO LIMIT RADIATIVE CORRECTIONS
TO THE COSMOLOGICAL CONSTANT BY M_{Susy}^4

Ulrich ELLWANGER

Laboratoire de Physique Théorique et Hautes Energies*

Université de Paris XI, bâtiment 211, 91405 Orsay Cedex, France

Abstract

Supergravity models are constructed in which the effective low energy theory contains only “super-soft” explicit supersymmetry breaking: masses of the scalars and pseudoscalars within a multiplet are split in opposite directions. With this form of supersymmetry breaking the radiative corrections of the matter sector to the vacuum energy are bounded by $\mathcal{O}(M_{\text{Susy}}^4)$ to all orders in perturbation theory, and we require $\text{Str } M^2 = 0$ including the hidden sector. The models are based on Kähler potentials obtained in recent orbifold compactifications, and we describe the construction of realistic theories.

LPTHE Orsay 94-106

December 1994

* Laboratoire associé au Centre National de la Recherche Scientifique - URA 63

The smallness of the cosmological constant belongs to the unsolved problems in particle physics. Within the standard model, we obtain a vacuum energy and hence a cosmological constant of $\mathcal{O}(M_{Weak}^4)$ [1].

The situation becomes even worse, if we extend the validity of the standard model (with the help of fine tuning in order to keep M_{Weak} stable) up to larger scales as M_{GUT} . At least quantum corrections then generate a vacuum energy of $\mathcal{O}(M_{GUT}^4)$. Supersymmetry, which helps to keep M_{Weak} stable under radiative corrections without fine tuning, is also of some help concerning the problem of the cosmological constant. Whereas unbroken global supersymmetry implies a vanishing vacuum energy [2], soft explicit susy breaking (of $\mathcal{O}(M_{Susy}) \sim \mathcal{O}(M_{Weak})$) leads to non-vanishing radiative corrections to the vacuum energy.

The magnitude of these contributions depends on the type of the soft susy breaking terms. Using M_{Planck} as an ultraviolet cutoff, the typically assumed “semi-soft” gaugino masses, trilinear couplings among the scalars and identical masses for the scalar and pseudoscalar components of a chiral multiplet generate a vacuum energy of $\mathcal{O}(M_{Susy}^2 \cdot M_{Planck}^2)$. There exists a further possibility to break supersymmetry explicitly, but softly: if we denote the complex scalar component of a chiral multiplet by A , a term of the form

$$\frac{m^2}{2}(A^2 + \overline{A}^2) \tag{1}$$

gives opposite contributions to the masses squared of its real and imaginary (or scalar and pseudoscalar) components. After expressing the explicit susy breaking in terms of a spurion field in superspace [3] and using the corresponding power counting rules it is straightforward to show, that in this case the maximal radiative contributions to the vacuum energy, to all orders in perturbation theory, are bounded by $\mathcal{O}(M_{Susy}^4)$ (with $m^2 \sim M_{Susy}^2$) [4]. This is certainly a big step forward compared with $\mathcal{O}(M_{Susy}^2 \cdot M_{Planck}^2)$. The generally assumed “semi-soft” form of explicit susy breaking thus does not exploit the potential power of supersymme-

try to bound the magnitude of radiative corrections in the same way as a cutoff $\Lambda \sim M_{Susy}$ would do.

Subsequently we will consider supergravity theories which are motivated by superstring models (typically involving orbifold compactifications). Recently problems related to the minimization of the potential in the presence of contributions to the vacuum energy of $\mathcal{O}(M_{Susy}^2 \cdot M_{Planck}^2)$ have been pointed out for this kind of theories [5].

To one loop level, the absence of radiative contributions of $\mathcal{O}(M_{Susy}^2 \cdot M_{Planck}^2)$ to the Coleman-Weinberg effective potential [6] boils down to the condition of a vanishing $Str\ M^2$. General supergravity theories with $Str\ M^2 = 0$ have been discussed in [7, 8]. In the presence of explicit supersymmetry breaking one needs, however, a guiding principle like super-soft susy breaking; otherwise, higher loop orders will inevitably generate contributions of $\mathcal{O}(M_{Susy}^2 \cdot M_{Planck}^2)$, even if they happen to be absent at the one loop level.

The questions arise, whether within string motivated theories such models with only super-soft susy breaking can be obtained, and whether these models can be realistic. Both questions will be answered positively below. Actually, supergravity theories with super-soft susy breaking have already been developed and discussed in [4, 9]. These theories were constructed such that the minimization of the vacuum energy generates automatically a hierarchy $M_{Susy} \ll M_{Planck}$ (which, for self-consistency, requires the vacuum energy to be of $\mathcal{O}(M_{Susy}^4)$). The resulting constraints on the parameters and particle spectrum are, however, very tight [10]. We will not discuss the generation of the hierarchy M_{Susy} vs. M_{Planck} here, but concentrate on general features of a vacuum energy of $\mathcal{O}(M_{Susy}^4)$.

The problem is complicated by the fact, that generally the light particle content of these theories is split into an “observable” matter sector and a “hidden” sector (typically the graviton, dilaton and moduli superfields), whose interactions with

the matter sector are suppressed by inverse powers of M_{Planck} . Only the matter sector makes up the softly broken supersymmetric theory discussed up to now, and higher loop corrections involving the hidden sector are beyond the scope of this paper. We will take the following attitude: due to the weak interactions between the matter and hidden sectors we require that each sector generates a maximal contribution to the vacuum energy of $\mathcal{O}(M_{Susy}^4)$ separately. To one loop level we are able to control the hidden sector with the help of the vanishing of $Str M^2$; within the matter sector with only super-soft supersymmetry breaking $Str M^2$ vanishes automatically [11]. To higher, and actually arbitrary loop level we are able to control the matter sector with the help of the above-mentioned theorem on the vacuum energy in super-softly broken susy. This way we will obtain realistic theories, whose behaviour towards the vacuum energy is as gentle as possible within nowadays available technologies.

In order to construct the supergravity theory we have to specify the Kähler potential K , the superpotential W and the gauge kinetic function f [12]. The chiral superfields within the matter sector will be denoted by the letters C_i and A_i ; the total number of C_i and A_i fields will be denoted by N_c . The hidden sector contains a dilaton S and N_M moduli fields denoted by T, M_i with $i = 1 \dots N_M - 1$. The moduli field T is singled out to play the role of an overall “breathing” mode. The total number of fields is thus given by $N = N_c + N_M + 1$.

The purpose of this paper is to propose a supergravity theory, which leads to an effective low energy theory with exclusively super-soft susy breaking. Its Kähler potential is given by

$$\begin{aligned}
K = & -\ell n(S + \bar{S}) - 3\ell n(T + \bar{T} - C_i \bar{C}_i - A_i \bar{A}_i) \\
& + h(T, \bar{T}) (A_i \bar{A}_i + \bar{A}_i \bar{A}_i) + \tilde{K}(M_i, \bar{M}_i) \quad .
\end{aligned}
\tag{2}$$

The first two terms are familiar from the construction of four-dimensional superstring theories [13], the second already from “No-scale” [14] or $SU(N, 1)$ [4, 9]

supergravity theories. The third term, involving the function $h(T, \bar{T})$, has been proposed as a solution of the so-called μ -problem of the $MSSM$ in [15]. Recently it has also been shown to appear in orbifold compactified superstring theories [16, 17]. In order to allow for the term of the form $A_i A_i + h.c.$, the fields A_i have to transform as real representations under all internal symmetries. (Of course, $A_i A_i + h.c.$ could be replaced by $A_i B_i$ with B_i transforming as the complex conjugate representation of A_i). The part of the Kähler potential involving the moduli M_i , $\tilde{K}(M_i, \bar{M}_i)$, is just required to lead to a positive definite metric $\tilde{K}^{M_i}_{\bar{M}_i}$.

For the superpotential W we make the ansatz

$$W = \mu(S, M_i) + \tilde{W}(C_i, A_i) \quad (3)$$

where every term in \tilde{W} is cubic in the fields C_i, A_i . The first term μ is a familiar result of gaugino condensation in the matter sector [18]. For the gauge kinetic function f we also assume a field dependence of the form

$$f(S, M_i) \quad . \quad (4)$$

Below we will assume a form of the function $h(T, \bar{T})$ such that all components of the fields A_i have positive masses squared and hence vanishing vevs. With $\langle A_i \rangle = 0$ one finds the identity ($\mathcal{G} = K + \ell n|W|^2$)

$$\mathcal{G}_{,T} (\mathcal{G}^{,T}_{\bar{T}})^{-1} \mathcal{G}^{,\bar{T}} = 3 \quad . \quad (5)$$

The tree level scalar potential [12]

$$V_{Tree} = e^{\mathcal{G}} \left(\mathcal{G}_{,I} (\mathcal{G}^{,I}_{\bar{J}})^{-1} \mathcal{G}^{,\bar{J}} - 3 \right) \quad (6)$$

is then easily seen to be positive semi-definite, and minimized for

$$\mathcal{G}_{,S} = \mathcal{G}_{,M_i} = \mathcal{G}_{,C_i} = \mathcal{G}_{,A_i} = 0 \quad . \quad (7)$$

Thus one obtains a “Goldstino angle” [19] $\theta = 0$. The vev of the field T is undetermined at this level, but in any case the tree level vacuum energy vanishes exactly. Supersymmetry is spontaneously broken; the corresponding gravitino mass is given by (in the units $M_{Planck}/8\pi = 1$)

$$m_{3/2} = e^{\mathcal{G}} = \frac{\mu e^{\tilde{K}}}{(S + \bar{S})(T + \bar{T})^3} \quad . \quad (8)$$

In order to derive the effective low energy theory for the matter sector it is of considerable help to note that among all possible \mathcal{G}_I only \mathcal{G}_T is nonzero. From $\mathcal{G}_S = \mathcal{G}_{M_i} = 0$ it follows, e.g., that no susy breaking gaugino masses are present at tree level. From an investigation of the scalar and fermionic interactions, and an appropriate rescaling of the fields in order to achieve canonical kinetic energies, one finds that the effective low energy theory is described by a superpotential W_{eff} of the form

$$W_{eff} = \frac{1}{(S + \bar{S})^{1/2}} \tilde{W}(C_i, A_i) + \frac{M_A}{2} A_i A_i \quad (9)$$

with

$$M_A = 2m_{3/2}(T + \bar{T}) (h + (T + \bar{T})h_{, \bar{T}}) \quad . \quad (10)$$

The only susy breaking interactions are indeed of the form of eq. (1), with

$$m^2 = -2m_{3/2}^2(T + \bar{T})^2 (h_{, T} + h_{, \bar{T}} + (T + \bar{T})h_{, T\bar{T}}) \quad . \quad (11)$$

Thus we have accomplished the first part of our task. Next we turn to the condition of the vanishing supertrace $Str M^2$, in order to tame one loop contributions to the vacuum energy of $\mathcal{O}(M_{Susy}^2 \cdot M_{Planck}^2)$ including the hidden sector. In the present case of vanishing gaugino masses, the formula of [20] (see also [8]) for $Str M^2$ for general supergravity theories boils down to

$$Str M^2 = 2m_{3/2}^2 \left[N - 1 - \mathcal{G}^{, I} \partial_I \partial_{\bar{J}} \ln \det(\mathcal{G}_{, M\bar{N}}) \mathcal{G}^{, \bar{J}} \right] \quad . \quad (12)$$

In our case, with $\mathcal{G}^{, I} = 0$ except for $I = T$, we only need

$$\mathcal{G}^{, T} \partial_T \partial_{\bar{T}} \ln \det(\mathcal{G}_{, M\bar{N}}) \mathcal{G}^{, \bar{T}} = 2 + N_c + \mathcal{O}(A, \bar{A}) \quad . \quad (13)$$

With $\langle A \rangle = 0$ and $N = N_c + N_M + 1$ we thus obtain

$$Str M^2 = 2m_{3/2}^2 [N_M - 2] \quad , \quad (14)$$

and the condition of a vanishing supertrace just becomes $N_M = 2$. This way we have constructed a model of the class discussed in [8]; we have obtained, however, a model with a special property: since, within the observable sector, susy is only broken by mass terms of the form of eq. (1), the part of the supertrace to which the observable sector contributes vanishes by itself [11]. Hence, with $N_M = 2$, the contribution of the hidden sector to the supertrace vanishes by itself as well. Thus we have satisfied all necessary conditions, which are required in order that the contributions to the vacuum energy are limited by $\mathcal{O}(M_{Susy}^4)$ from the observable sector to all orders in perturbation theory, and those from the hidden sector to one loop order.

Let us now discuss the possibility of constructing realistic models within this class of theories. Apart from a reasonable particle content we require sufficiently large masses for the gauginos, squarks and sleptons, and negative masses squared in the Higgs sector in order to trigger $SU(2) \times U(1)$ symmetry breaking. Generally this is achieved, if not already at tree level, with the help of radiative corrections. In the presence of semi-hard soft terms, the most important radiative corrections are logarithmically divergent and most conveniently summed up by integrating the renormalization group equations from the cutoff scale M_{Plank} down to the weak scale or M_{Susy} .

In a model with exclusive super-soft susy breaking, however, the only logarithmically divergent contributions to the effective action (apart from wave function normalizations) are proportional to F -components of singlet superfields [3]. Consequently, the right hand sides of the RG equations for gaugino, squark and slepton masses as well as trilinear scalar couplings vanish. Nevertheless such masses are radiatively generated; now, however, the corresponding Feynman diagrams are

ultraviolet finite.

Let us first have a look at the radiative corrections involving gauge interactions. We assume that the fields A_i , with masses M_A and mass splittings as in eq. (1), transform as representations r under gauge groups a . For $m^2 \ll M_A^2$ the corresponding gauginos receive one loop masses given by

$$m_a = \sum_r \frac{\alpha_a m^2}{4\pi M_A} T_r^a \quad . \quad (15)$$

The Casimir eigenvalue T_r^a has to be replaced by the charge squared in the case of a $U(1)$ gauge group, and is given by $1/2$ resp. N for r denoting a fundamental resp. adjoint representation of $SU(N)$. At two loop order, all scalars of the theory which transform as representations q under the gauge groups a obtain positive definite masses squared given by [21]

$$m_q^2 = \sum_{a,r} \frac{\alpha_a^2 m^4}{8\pi^2 M_A^2} T_r^a C_q^a \quad . \quad (16)$$

In the case of $U(1)$ C_q^a is again given by the square of the charge of q , whereas $C_q^a = \frac{N^2-1}{2N}$ resp. N for fundamental resp. adjoint representations of $SU(N)$.

The most important Yukawa mediated radiative corrections are the above-mentioned logarithmically divergent ones, if a singlet field couples to the massive fields A . Let us assume a corresponding term in the superpotential W_{eff} ,

$$W_{eff} = \beta S A A + W' \quad . \quad (17)$$

Then, to one loop order, the following contribution to the effective potential is obtained:

$$-\frac{\beta}{16\pi^2} m^2 (F_S + \overline{F}_S) \ell n \frac{\Lambda^2}{M_A^2} \quad (18)$$

where F_S is an expression quadratic in the scalar fields given by $F_S = \frac{\partial W_{eff}}{\partial S}$. The term (18) is actually the presently harmless result of the singlet tadpole contribution, which can mess up hierarchies in the presence of semi-soft susy breakings

[22]. Further Yukawa induced radiative corrections are generally negligible for $m^2 \ll M_A^2$; the interaction in eq. (17), e.g., gives additionally rise to a positive mass m_S^2 for S of

$$m_S^2 = \frac{\beta^2}{8\pi^2} \frac{m^6}{M_A^4} \quad . \quad (19)$$

and a term linear in S

$$\frac{\beta}{16\pi^2} \frac{m^4}{M_A} (S + \bar{S}) \quad . \quad (20)$$

These results can already be used to discuss the construction of realistic models. First, in order to generate realistic gaugino masses for all gauginos of the standard model gauge groups via (15), the fields A_i should carry quantum numbers under all these gauge groups. Second, in order for gaugino, squark and slepton masses via (15) and (16) to be sufficiently large, the ratio $\frac{m^2}{M_A}$ should be at least of $\mathcal{O}(TeV)$. Both arguments rule out the attractive possibility of identifying AA with the *MSSM* Higgs fields $H_1 H_2$ as in [23] (and choosing $m^2 > M_A^2$ in order to generate Higgs VEVs already at tree level).

Instead, the fields A have to be identified with new fields beyond the *MSSM*, with $M_A > \mathcal{O}(TeV)$. The successful unification of the running gauge couplings at M_{GUT} within the *MSSM* is not spoiled, if the fields A are chosen to fill up complete representations of $SU(5)$ as, e.g., $5 + \bar{5}$.

A realistic class of models can thus be constructed as follows: we identify the fields C_i of eqs. (2) and (3) with the quark, lepton and Higgs superfields of the *MSSM* as well as a gauge singlet superfield S . In addition the matter sector contains the fields A_i . For the cubic part \tilde{W} of the superpotential in eq. (3) we choose

$$\tilde{W} = \tilde{\lambda} S H_1 H_2 + \frac{\tilde{k}}{3} S^3 + \tilde{\beta} S A_i A_i + \dots \quad (21)$$

where the dots denote the *MSSM* Yukawa couplings. According to eq. (9) the effective superpotential W_{eff} then reads

$$W_{eff} = \lambda S H_1 H_2 + \frac{k}{3} S^3 + \beta S A_i A_i + \frac{M_A}{2} A_i A_i + \dots \quad (22)$$

with rescaled Yukawa couplings and M_A given by eq. (10). Note that for generic functions $h(T, \bar{T})$ and vevs of T of $\mathcal{O}(M_{Planck})$ we have from eq. (11)

$$M_A \sim m \sim M_{Susy} \quad ; \quad (23)$$

the condition $M_A > m$ turns into a condition on the form of $h(\bar{T}, T)$ and the vev of T . ($M_A \gg m$ just simplifies the computation of the radiative corrections.)

Now, at one resp. two loop order, gaugino as well as positive squark, slepton and Higgs masses are generated according to eqs. (15) and (16). The logarithmically divergent contribution (18), with $F_S = \lambda H_1 H_2 + k S^2 + \beta A_i A_i$, generates Higgs masses of the form

$$m_3^2 H_1 H_2 + h.c. \quad (24a)$$

with

$$m_3^2 = -\frac{\lambda \beta d_r}{16\pi^2} m^2 \ell n \frac{\Lambda^2}{M_A^2} \quad , \quad (24b)$$

which can destabilize the Higgs potential as desired. (d_r denotes the dimension of the representation r of the fields A_i). Actually, the positive two loop stop masses m_{st}^2 from eq. (16) induce, at three loop order, negative masses squared for the Higgs fields H_2 , which have Yukawa couplings h_t to the top quarks [24]. For $m_{st}^2 \ll m^2$ this term can be estimated to be

$$-\frac{3}{4\pi^2} h_t^2 m_{st}^2 |H_2|^2 \ell n \frac{M_A}{m_{st}} \quad . \quad (25)$$

With m_{st}^2 given by eq. (16), where α_{QCD} gives the leading contribution, this negative mass for H_2 can be numerically larger than the positive mass for H_2 obtained via eq. (16), where only the electroweak gauge couplings appear. Thus there exist even two possible mechanisms to trigger the desired $SU(2) \times U(1)$ symmetry breaking.

We have checked that, after minimization of the complete scalar potential including the S dependent terms (19) and (20), a wide range of parameters λ, k, β, M_A

and m exists, which leads to realistic particle masses. The details of the spectrum depend, in addition, on the representation r of the fields A_i . A complete analysis of the full range of parameters is beyond the scope of the present paper. Some limiting situations are, however, worth mentioning: one may choose the Yukawa coupling k vanishingly small or even equal to zero. Then it is only the mass m_S^2 of eq. (19), which stabilizes the potential for S , and S assumes a large vev of $\mathcal{O}(M_A^3/\beta m^2)$. Accordingly the ratio M_A/m has not to be chosen too large, and λ has to be small in order to allow for vevs of H_1 and H_2 . The spectrum will contain a light pseudoscalar (dominantly gauge singlet), since the Peccei-Quinn $U(1)$ symmetry in the H_1, H_2, S - sector is only weakly broken by the terms (18), (20).

Independently thereof, if the F term (24) plays the dominant role in triggering nonvanishing vevs of H_1 and H_2 , $\tan \beta = \frac{\langle H_2 \rangle}{\langle H_1 \rangle}$ will be close to 1, whereas $\tan \beta$ will be large in the case where the h_t -induced contribution (25) is most important.

Let us conclude with some comments on the prospects of realizing the ansatz for K (eq. (2)), W (eq. (3)) and f (eq. (4)) within string theory. The first two terms in eq. (2) and the form of f in eq. (4) are well known from string theory at tree level. Also a function $h(T, \bar{T})$ has been found to arise at tree level [16, 17]; the corresponding expression is of the form

$$h(T, \bar{T}) \sim \frac{1}{T + \bar{T}} \quad . \quad (26)$$

With such a function $h(T, \bar{T})$ one finds, from eqs. (10) and (11), $M_A = m = 0$. However, string loop corrections generate functions $h(T, \bar{T})$ different from (26) [17]. Generally they also modify the first two terms of eq. (2) [25], and the ansatz for f eq. (4) [26]. The details depend, of course, on the string model under consideration (as, e.g., on the Green Schwarz angles δ_{GS}^i). For the present class of models to emerge it would thus be desirable that string loop corrections leave the Kähler potential K and the gauge kinetic function f untouched, but modify the

form of the function $h(T, \overline{T})$. It remains to be seen, however, whether such string models can be constructed.

The purpose of this paper is to point out the existence of supergravity theories with exclusive supersoft supersymmetry breaking, in which even radiative corrections generate a vacuum energy of not more than $\mathcal{O}(M_{Susy}^4)$. We emphasize the possibility of constructing realistic models within these theories along the lines discussed above. One could imagine that string theory would like to make use of this attractive possibility to control the quantum corrections to the cosmological constant. This would lead to important restrictions on the string models resp. string vacua.

References

- [1] A. Linde, JETP Letters 19 (1974) 183.
- [2] B. Zumino, Nucl. Phys. B89 (1975) 535.
- [3] L. Girardello, M. Grisaru, Nucl. Phys. B194 (1982) 65.
- [4] U. Ellwanger, N. Dragon, M. Schmidt, Nucl. Phys. B255 (1985) 549.
- [5] K. Choi, J. E. Kim, H. P. Nilles, Phys. Rev. Lett. 73 (1994) 1758.
- [6] S. Coleman, E. Weinberg, Phys. Rev. D7 (1973) 1888.
- [7] S. Ferrara, C. Kounnas, M. Porrati, F. Zwirner, Phys. Lett. 194B (1987) 366,
C. Kounnas, M. Quiros, F. Zwirner, Nucl. Phys. B302 (1988) 403.
- [8] S. Ferrara, C. Kounnas, F. Zwirner, CERN preprint TH-7192/94.
- [9] N. Dragon, U. Ellwanger, M. Schmidt, Phys. Lett. 145B (1984) 192, Phys.
Lett. 154B (1985) 373, Z. Phys. C29 (1985) 219, Progr. in Part. and Nucl.
Phys. 18 (1987) 1.
- [10] U. Ellwanger, Z. Phys. C47 (1990) 281.
- [11] S. Ferrara, L. Girardello, F. Palumbo, Phys. Rev. D20 (1979) 403.
- [12] E. Cremmer, S. Ferrara, L. Girardello, A. Van Proyen, Nucl. Phys. B212 (1983)
413.
- [13] E. Witten, Phys. Lett. B155 (1985) 151.
- [14] E. Cremmer, S. Ferrara, C. Kounnas, D. Nanopoulos, Phys. Lett. B133 (1983)
61; A. Lahanas, D. Nanopoulos, Phys. Rep. 145 (1987) 1.
- [15] G. Giudice and A. Masiero, Phys. Lett. B206 (1988) 480.
- [16] G. Lopes Cardoso, D. Lüst, T. Mohaupt, Nucl. Phys. B432 (1994) 68.
- [17] I. Antoniadis, E. Gava, K. Narain, T. Taylor, Nucl. Phys. B432 (1994) 187.
- [18] J.-P. Derendinger, L. Ibáñez, H.P. Nilles, Phys. Lett. B155 (1985) 65;
M. Dine, R. Rohm, N. Seiberg, E. Witten, Phys. Lett. B156 (1985) 55.
- [19] A. Brignole, L. Ibáñez, C. Muñoz, Nucl. Phys. B422 (1994) 125.
- [20] M. Grisaru, M. Rocek, A. Karlhede, Phys. Lett. B120 (1983) 110.

- [21] M. Dine, W. Fischler, Nucl. Phys. B204 (1982) 346;
L. Alvarez-Gaumé, M. Claudson, M. Wise, Nucl. Phys. B207 (1982) 96.
- [22] J. Polchinski, L. Susskind, Phys. Rev. D26 (1982) 3661;
H.-P. Nilles, M. Srednicki, D. Wyler, Phys. Lett. B124 (1983) 337;
A. Lahanas, Phys. Lett. B124 (1983) 341;
U. Ellwanger, Phys. Lett. B133 (1983) 187;
J. Bagger and E. Poppitz, Phys. Rev. Lett. 71 (1993) 2380.
- [23] A. Brignole, F. Zwirner, preprint CERN-TH.7439/94, hep-th 9409099.
- [24] L. Ibáñez, G. G. Ross, Phys. Lett. 110B (1982) 215.
- [25] J.-P. Derendinger, S. Ferrara, C. Kounnas, F. Zwirner, Phys. Lett. B271 (1991) 307 and Nucl. Phys. B372 (1992) 145;
I. Antoniadis, E. Gava, K. Narain, T. Taylor, Nucl. Phys. B407 (1993) 706.
- [26] L. Ibáñez, H.-P. Nilles, Phys. Lett. B169 (1986) 354;
H. Itoyama, J. Leon, Phys. Rev. Lett. 56 (1986) 2352;
L. Dixon, V. Kaplunovsky, J. Louis, Nucl. Phys. B255 (1991) 649.